

INDUCED-GRAVITY INFLATION IN NO-SCALE SUPERGRAVITY AND BEYOND

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ABSTRACT

Supersymmetric versions of induced-gravity inflation are formulated within Supergravity (SUGRA) employing two gauge singlet chiral superfields. The proposed superpotential is uniquely determined by applying a continuous R and a discrete \mathbb{Z}_n symmetry. We select two types of logarithmic Kähler potentials, one associated with a no-scale-type $SU(2, 1)/SU(2) \times U(1)_R \times \mathbb{Z}_n$ Kähler manifold and one more generic. In both cases, imposing a lower bound on the parameter $c_{\mathcal{R}}$ involved in the coupling between the inflaton and the Ricci scalar curvature – e.g. $c_{\mathcal{R}} \gtrsim 76, 105, 310$ for $n = 2, 3$ and 6 respectively –, inflation can be attained even for subplanckian values of the inflaton while the corresponding effective theory respects the perturbative unitarity. In the case of no-scale SUGRA we show that, for every n , the inflationary observables remain unchanged and in agreement with the current data while the inflaton mass is predicted to be $3 \cdot 10^{13}$ GeV. Beyond no-scale SUGRA the inflationary observables depend mildly on n and crucially on the coefficient involved in the fourth order term of the Kähler potential which mixes the inflaton with the accompanying non-inflaton field.

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1 INTRODUCTION

The announcement of the recent PLANCK results [1, 2] fuelled increasing interest in inflationary models implemented thanks to a strong enough non-minimal coupling between the inflaton field, ϕ , and the Ricci scalar curvature, \mathcal{R} . Indeed, these models predict [2, 3] a (scalar) spectral index n_s , tantalizingly close to the value favored by observational data. The existing non-minimally coupled to Gravity inflationary models can be classified into two categories depending whether the non-minimal coupling to \mathcal{R} is added into the conventional one, $m_{\text{P}}^2 \mathcal{R}/2$ – where $m_{\text{P}} = 2.44 \cdot 10^{18}$ GeV is the reduced Planck scale – or it replaces the latter. In the first case the *vacuum expectation value* (v.e.v) of the inflaton after inflation assumes sufficiently low values after inflation, such that a transition to Einstein gravity at low energy to be guaranteed. In the second case, however, the term $m_{\text{P}}^2 \mathcal{R}/2$ is dynamically generated via the v.e.v of the inflaton; these models are, thus, named [4, 5] *Induced-Gravity* (IG) inflationary models. Despite the fact that both models of *non-Minimal Inflation* (nMI) are quite similar during inflation and may be collectively classified into universal “attractor” models [6], they exhibit two crucial differences. Namely, in the second category, (i) the *Einstein frame* (EF) inflationary potential develops a singularity at $\phi = 0$ and so, inflation is of Starobinsky-type [7] actually; (ii) The *ultraviolet* (UV) cut-off scale [8–10] of the theory, as it is recently realized [11, 12], can be identified with m_{P} and, thereby, concerns regarding the naturalness of inflation can be safely eluded. On the other hand, only some [10] of the remaining models of nMI can be characterized as unitarity safe.

In a recent paper [11] a *supersymmetric* (SUSY) version of IG inflation was, for first time, presented within no-scale [13–15] *Supergravity* (SUGRA). A Higgs-like modulus plays there the role of inflaton, in sharp contrast to Ref. [14] where the inflaton is matter-like. For this reason we call in Ref. [11] the inflationary model *no-scale modular inflation*. Although any connection with the no-scale SUSY breaking [13, 16] is lost in that setting, we show that the model provides a robust cosmological scenario linking together non-thermal leptogenesis, neutrino physics and a resolution to the μ problem of the *Minimal SUSY SM* (MSSM). Namely, in Ref. [11], we employ a Kähler potential, K , corresponding to a $SU(N, 1)/SU(N) \times U(1)_R \times \mathbb{Z}_2$ symmetric Kähler manifold. This symmetry fixes beautifully the form of K up to an holomorphic function Ω_{H} which exclusively depends on the inflaton, ϕ , and its form $\Omega_{\text{H}} \sim \phi^2$ is fixed by imposing a \mathbb{Z}_2 discrete symmetry which is also respected by the superpotential W . Moreover, the model possesses a continuous R symmetry, which reduces to the well-known R -parity

of MSSM. Thanks to the strong enough coupling between ϕ and \mathcal{R} , inflation can be attained even for subplanckian values of ϕ , contrary to other SUSY realizations [15, 17, 18] of the Starobinsky-type inflation.

Most recently a more generic form of Ω_H has been proposed [12] at the non-SUSY level. In particular, Ω_H is specified as $\Omega_H \sim \phi^n$ and it was pointed out that the resulting IG inflationary models exhibit an attractor behavior since the inflationary observables and the mass of the inflaton at the vacuum are independent of the choice of n . It would be, thereby, interesting to investigate if this nice feature insists also in the SUSY realizations of these models. This aim gives us the opportunity to generalize our previous analysis [11] and investigate the inflationary predictions independently of the post-inflationary cosmological evolution. Namely, we here impose on Ω_H a discrete \mathbb{Z}_n symmetry with $n \geq 2$, and investigate its possible embedding in standard Poincaré SUGRA, without invoking the superconformal formulation – cf. Ref. [19]. We discriminate two possible embeddings, one based on a no-scale-type symmetry and one more generic, with the first of these being much more predictive. Namely, while the embedding of IG models in generic SUGRA gives adjustable results as regards the inflationary observables, – see also Ref. [20] –, no-scale SUGRA predicts independently of n results identical to those obtained in the non-SUSY case. Therefore, no-scale SUGRA consists a natural framework in which such models can be implemented.

Below, in Sec. 2, we describe the generic formulation of IG models within SUGRA. In Sec. 3 we present the basic ingredients of our IG inflationary models, derive the inflationary observables and confront them with observations. We also provide a detailed analysis of the UV behavior of these models in Sec. 4. Our conclusions are summarized in Sec. 5. Throughout the text, the subscript of type χ denotes derivation *with respect to* (w.r.t) the field χ (e.g., $_{,\chi\chi} = \partial^2/\partial\chi^2$) and charge conjugation is denoted by a star.

2 EMBEDDING IG INFLATION IN SUGRA

In Sec. 2.1 we present the basic formulation of a theory which exhibits non-minimal coupling of scalar fields to \mathcal{R} within SUGRA and in Sec. 2.2 we outline our strategy in constructing viable models of IG inflation. The general framework for the analysis of the emerged models is given in Sec. 2.3.

2.1 THE GENERAL SET-UP

Our starting point is the EF action for N gauge singlet scalar fields z^α within SUGRA [21, 22] which can be written as

$$S = \int d^4x \sqrt{-\hat{\mathbf{g}}} \left(-\frac{1}{2} m_{\text{P}}^2 \hat{\mathcal{R}} + K_{\alpha\bar{\beta}} \hat{g}^{\mu\nu} \partial_\mu z^\alpha \partial_\nu z^{*\bar{\beta}} - \hat{V} \right), \quad (2.1a)$$

where summation is taken over the scalar fields z^α , $K_{\alpha\bar{\beta}} = \hat{K}_{,z^\alpha z^{*\bar{\beta}}}$ with $K^{\bar{\beta}\alpha} K_{\alpha\bar{\gamma}} = \delta_{\bar{\gamma}}^{\bar{\beta}}$, $\hat{\mathbf{g}}$ is the determinant of the EF metric $\hat{g}_{\mu\nu}$, $\hat{\mathcal{R}}$ is the EF Ricci scalar curvature, \hat{V} is the EF F-term SUGRA scalar potential which can be extracted once the superpotential W and the Kähler potential K have been selected, by applying the standard formula

$$\hat{V} = e^{K/m_{\text{P}}^2} \left(K^{\alpha\bar{\beta}} F_\alpha F_{\bar{\beta}}^* - 3 \frac{|W|^2}{m_{\text{P}}^2} \right), \quad \text{where } F_\alpha = W_{,z^\alpha} + K_{,z^\alpha} W / m_{\text{P}}^2. \quad (2.1b)$$

Note that D-term contributions into \hat{V} do not exist since we consider gauge singlet z^α 's. By performing a conformal transformation and adopting a frame function Ω which is related to K as follows

$$-\Omega/3 = e^{-K/3m_{\text{P}}^2} \Rightarrow K = -3m_{\text{P}}^2 \ln(-\Omega/3), \quad (2.2)$$

we arrive at the following action

$$S = \int d^4x \sqrt{-g} \left(-\frac{m_P^2}{2} \left(-\frac{\Omega}{3} \right) \mathcal{R} + m_P^2 \Omega_{\alpha\bar{\beta}} \partial_\mu z^\alpha \partial^\mu z^{*\bar{\beta}} - \Omega \mathcal{A}_\mu \mathcal{A}^\mu / m_P^2 - V \right), \quad (2.3)$$

where $g_{\mu\nu} = -(3/\Omega) \hat{g}_{\mu\nu}$ and $V = \Omega^2 \hat{V}/9$ are the JF metric and potential respectively, we use the shorthand notation $\Omega_\alpha = \Omega_{,z^\alpha}$ and $\Omega_{\bar{\alpha}} = \Omega_{,z^{*\bar{\alpha}}}$ and \mathcal{A}_μ is the purely bosonic part of the on-shell value of an auxiliary field given by

$$\mathcal{A}_\mu = -im_P^2 (\Omega_\alpha \partial_\mu z^\alpha - \Omega_{\bar{\alpha}} \partial_\mu z^{*\bar{\alpha}}) / 2\Omega. \quad (2.4)$$

It is clear from Eq. (2.3) that S exhibits non-minimal couplings of the z^α 's to \mathcal{R} . However, Ω enters the kinetic terms of the z^α 's too. In general, Ω can be written as [21]

$$-\Omega/3 = \Omega_H(z^\alpha) + \Omega_H^*(z^{*\bar{\alpha}}) - \Omega_K(z^\alpha z^{*\bar{\alpha}})/3, \quad (2.5)$$

where Ω_K is a dimensionless real function while Ω_H is a dimensionless, holomorphic function. For $\Omega_H > \Omega_K$, Ω_K expresses mainly the kinetic terms of the z^α 's whereas Ω_H represents the non-minimal coupling to gravity – note that $\Omega_{\alpha\bar{\beta}}$ is independent of Ω_H since $\Omega_{H,z^\alpha z^{*\bar{\beta}}} = 0$.

To realize the idea of IG, we have to assume that Ω_H depends on a Higgs-like modulus, $z^1 := \Phi$ whose the v.e.v generates the conventional term of the Einstein gravity at the SUSY vacuum, i.e.

$$\langle \Omega_H \rangle + \langle \Omega_H^* \rangle = 1 \Rightarrow \langle \Omega_H \rangle = 1/2 \text{ for } \langle \Omega_K \rangle \sim 0 \quad (2.6)$$

where we take into account that the phase of Φ , $\arg \Phi$ is stabilized to zero; we thus get $\langle \Omega_H \rangle = \langle \Omega_H^* \rangle$.

In order to get canonical kinetic terms, we need [21] $\mathcal{A}_\mu = 0$ and $\Omega_{K\alpha\bar{\beta}} \simeq 0$ or $\delta_{\alpha\bar{\beta}}$. The first condition is attained when the dynamics of the z^α 's is dominated only by the real moduli $|z^\alpha|$. The second condition is satisfied by the choice

$$\Omega_K(|z^\alpha|^2) = k_\alpha |z^\alpha|^2 / m_P^2 - k_{\alpha\beta} |z^\alpha|^2 |z^\beta|^2 / m_P^4 \quad (2.7)$$

with sufficiently small coefficients k_α and $k_{\alpha\beta} \simeq 1$. Here we assume that the z^α 's are charged under a global symmetry, so as mixed terms of the form $z^\alpha z_\beta^*$ are disallowed. The inclusion of the fourth order term for the accompanying non-inflaton field, $z^2 := S$ is obligatory in order to evade [21] a tachyonic instability occurring along this direction during IG inflation. As a consequence, all the allowed terms are to be considered in the analysis for consistency. Let us here note that such a consistency is not observed in the SUGRA incarnations of similar models [6, 21]. On the other hand, if we assume that

$$k_1 = 0 \text{ and } k_{1\alpha} = 0, \quad \forall \alpha = 1, \dots, N-1 \quad (2.8)$$

the emergent Kähler manifold associated with K can be identified with $SU(N, 1)/SU(N) \times U(1)_R \times \mathbb{Z}_n$ – where the symmetries $U(1)_R$ and \mathbb{Z}_n are specified in Sec. 2.2 – and highly simplifies the realization of IG inflation. The option in Eq. (2.8) is inspired by the early models of soft SUSY breaking [13] and defines [15] no-scale SUGRA. We below show details of these two realizations of IG inflation.

2.2 MODELING IG INFLATION IN SUGRA

As we anticipated above, the realization of the idea of IG in SUGRA requires at least two singlet superfields, i.e., $z^\alpha = S, \Phi$; Φ is a Higgs-like superfield whose the v.e.v generates m_P and S is an accompanying superfield, whose the stabilization at the origin assists us to isolate the contribution of Φ into \hat{V} , Eq. (2.1b). To see how this structure works, let us below specify the form of Ω_H and W .

Inspired by Ref. [12], we here determine Ω_H by postulating its invariance under the action of a global \mathbb{Z}_n discrete symmetry. Therefore it can be written as

$$\Omega_H(\Phi) = c_{\mathcal{R}} \frac{\Phi^n}{m_{\text{P}}^n} + \sum_{k=1}^{\infty} \lambda_k \frac{\Phi^{2kn}}{m_{\text{P}}^{2kn}} \quad (2.9)$$

with k being a positive integer. Restricting ourselves to subplanckian values of Φ and assuming relatively low λ_k 's, we can say that \mathbb{Z}_n uniquely determines the form of Ω_H . Confining ourselves to a such situation we ignore henceforth the k -dependent terms in Eq. (2.9). On the other hand, W has to be selected so as to achieve the arrangement of Eq. (2.6). The simplest choice is that used in the models of F-term hybrid inflation [23]. As a consequence $\Omega_H(\Phi)$ has to be involved also in the superpotential W of our model which has the form

$$W = \lambda m_{\text{P}}^2 S (\Omega_H - 1/2) / c_{\mathcal{R}} \quad (2.10)$$

and can be uniquely determined if we impose, besides \mathbb{Z}_n , a nonanomalous R symmetry $U(1)_R$ under which

$$S \rightarrow e^{i\alpha} S, \quad \Omega_H \rightarrow \Omega_H, \quad W \rightarrow e^{i\alpha} W. \quad (2.11)$$

Indeed, $U(1)_R$ symmetry ensures the linearity of W w.r.t S which is crucial for the success of our construction. To verify that W leads to the desired $\langle \Omega_H \rangle$ we minimize the SUSY limit, V_{SUSY} , of \hat{V} , obtained from the latter, when m_{P} tends to infinity. This is

$$V_{\text{SUSY}} = \lambda^2 m_{\text{P}}^4 |\Omega_H - 1/2|^2 / c_{\mathcal{R}}^2 + \lambda^2 m_{\text{P}}^4 |S \Omega_{H,\Phi}|^2 / c_{\mathcal{R}}^2, \quad (2.12a)$$

where the complex scalar components of Φ and S are denoted by the same symbol. From Eq. (2.12a), we find that the SUSY vacuum lies at

$$\langle S \rangle = 0 \quad \text{and} \quad \langle \Omega_H \rangle = 1/2, \quad (2.12b)$$

as required by Eq. (2.6). Let us emphasize that soft SUSY breaking effects explicitly break $U(1)_R$ to a discrete subgroup. Usually [11] combining the latter with the \mathbb{Z}_2^f fermion parity, yields the well-known R -parity of MSSM, which guarantees the stability of the lightest SUSY particle and therefore it provides a well-motivated CDM candidate.

The selected W and K by construction give also rise to a stage of IG inflation. Indeed, placing S at the origin, the only surviving term of \hat{V} in Eq. (2.1b) is

$$\hat{V}_{\text{IG0}} = e^{K/m_{\text{P}}^2} K^{SS^*} |W_{,S}|^2 = \frac{\lambda^2 m_{\text{P}}^4 |2\Omega_H - 1|^2}{4c_{\mathcal{R}}^2 f_{S\Phi} f_{\mathcal{R}}^2} \quad \text{since} \quad e^{K/m_{\text{P}}^2} = \frac{1}{f_{\mathcal{R}}^3} \quad \text{and} \quad K^{SS^*} = \frac{f_{\mathcal{R}}}{f_{S\Phi}}, \quad (2.13a)$$

where the functions $f_{\mathcal{R}}$ and $f_{S\Phi}$ are computed along the inflationary track, i.e.,

$$f_{\mathcal{R}} = -\Omega/3 \quad \text{and} \quad f_{S\Phi} = m_{\text{P}}^2 \Omega_{,SS^*} \quad \text{for} \quad S = \arg \Phi = 0. \quad (2.13b)$$

Given that $f_{S\Phi} \ll f_{\mathcal{R}} \simeq 2\Omega_H$ with $c_{\mathcal{R}} \gg 1$, an inflationary plateau emerges since the resulting \hat{V}_{IG0} in Eq. (2.13a) is almost constant. Therefore, Φ involved in the definition of Ω_H , Eq. (2.9), arises naturally as an inflaton candidate. Note that the non-vanishing values of Φ during IG inflation break spontaneously the imposed \mathbb{Z}_n ; no domain walls are thus produced due to the spontaneous breaking of \mathbb{Z}_n at the SUSY vacuum, Eq. (2.12b).

2.3 FRAMEWORK OF INFLATIONARY ANALYSIS

To consolidate the validity of the inflationary proposal we have to check the stability of the inflationary direction

$$\theta = s = \bar{s} = 0, \quad (2.14)$$

w.r.t the fluctuations of the various fields, which are expanded in real and imaginary parts as follows

$$\Phi = \frac{\phi}{\sqrt{2}} e^{i\theta/m_P} \quad \text{and} \quad S = \frac{s + i\bar{s}}{\sqrt{2}}. \quad (2.15)$$

To this end we examine the validity of the extremum and minimum conditions, i.e.,

$$\left. \frac{\partial \widehat{V}_{\text{IG}0}}{\partial \widehat{\chi}^\alpha} \right|_{\text{Eq. (2.14)}} = 0 \quad \text{and} \quad \widehat{m}_{\chi^\alpha}^2 > 0 \quad \text{with} \quad \chi^\alpha = \theta, s, \bar{s}. \quad (2.16a)$$

Here $\widehat{m}_{\chi^\alpha}^2$ are the eigenvalues of the mass matrix with elements

$$\widehat{M}_{\alpha\beta}^2 = \left. \frac{\partial^2 \widehat{V}_{\text{IG}0}}{\partial \widehat{\chi}^\alpha \partial \widehat{\chi}^\beta} \right|_{\text{Eq. (2.14)}} \quad \text{with} \quad \chi^\alpha = \theta, s, \bar{s} \quad (2.16b)$$

and hat denotes the EF canonically normalized fields. The kinetic terms of the various scalars in Eq. (2.1a) can be brought into the following form

$$K_{\alpha\beta} \dot{z}^\alpha \dot{z}^{\beta*} = \frac{1}{2} \left(\dot{\widehat{\phi}}^2 + \dot{\widehat{\theta}}^2 \right) + \frac{1}{2} \left(\dot{\widehat{s}}^2 + \dot{\widehat{\bar{s}}}^2 \right), \quad (2.17a)$$

where the dot denotes derivation w.r.t the JF cosmic time and the hatted fields are defined as follows

$$\frac{d\widehat{\phi}}{d\phi} = J = \sqrt{K_{\Phi\Phi*}}, \quad \widehat{\theta} = m_P \sqrt{K_{\Phi\Phi*}} \theta / \phi, \quad \text{and} \quad (\widehat{s}, \widehat{\bar{s}}) = \sqrt{K_{SS*}}(s, \bar{s}). \quad (2.17b)$$

Note, in passing, that the spinors ψ_Φ and ψ_S associated with the superfields S and Φ are normalized similarly, i.e., $\widehat{\psi}_S = \sqrt{K_{SS*}}\psi_S$ and $\widehat{\psi}_\Phi = \sqrt{K_{\Phi\Phi*}}\psi_\Phi$.

Upon diagonalization of $\widehat{M}_{\alpha\beta}^2$, Eq. (2.16b), we can construct the scalar mass spectrum of the theory along the direction in Eq. (2.14) – see Sec. 3.2.1 and 3.3.1. Besides the stability requirement in Eq. (2.16a), from the derived spectrum we can numerically verify that the various masses remain greater than \widehat{H}_{IG} during the last 50 e-foldings of inflation, and so any inflationary perturbations of the fields other than the inflaton are safely eliminated. Due to the large effective masses that θ, s and \bar{s} in Eq. (2.16b) acquire during inflation, they enter a phase of oscillations about zero with reducing amplitude. As a consequence, the ϕ dependence in their normalization – see Eq. (2.17b) – does not affect their dynamics. Moreover, we can observe that the fermionic (4) and bosonic (4) degrees of freedom are equal – here we take into account that $\widehat{\phi}$ is not perturbed. Employing the well-known Coleman-Weinberg formula [24], we find that the one-loop corrected inflationary potential is

$$\widehat{V}_{\text{IG}} = \widehat{V}_{\text{IG}0} + \frac{1}{64\pi^2} \left(\widehat{m}_\theta^4 \ln \frac{\widehat{m}_\theta^2}{\Lambda^2} + 2\widehat{m}_s^4 \ln \frac{\widehat{m}_s^2}{\Lambda^2} - 4\widehat{m}_{\psi_\pm}^4 \ln \frac{m_{\psi_\pm}^2}{\Lambda^2} \right), \quad (2.18)$$

where Λ is a renormalization group mass scale, \widehat{m}_θ and $\widehat{m}_s = \widehat{m}_{\bar{s}}$ are defined in Eq. (2.16a) and \widehat{m}_{ψ_\pm} are the mass eigenvalues which correspond to eigenstates $\widehat{\psi}_\pm \simeq (\widehat{\psi}_S \pm \widehat{\psi}_\Phi)/\sqrt{2}$. As we numerically verify, the one-loop corrections have no impact on our results, since the slope of the inflationary path is generated at the classical level and the various masses are proportional to the weak coupling λ .

3 THE INFLATIONARY SCENARIO

In this section we outline the salient features and the predictions of our inflationary scenaria in Secs. 3.2 and 3.3 respectively, testing them against a number of criteria introduced in Sec. 3.1.

3.1 INFLATIONARY OBSERVABLES – CONSTRAINTS

A successful inflationary scenario has to be compatible with a number of observational requirements which are outlined in the following.

3.1.1. The number of e-folds, \hat{N}_* , that the scale $k_* = 0.05/\text{Mpc}$ suffers during IG inflation,

$$\hat{N}_* = \int_{\hat{\phi}_f}^{\hat{\phi}_*} \frac{d\hat{\phi}}{m_{\text{P}}^2} \frac{\hat{V}_{\text{IG}}}{\hat{V}_{\text{IG},\hat{\phi}}} = \int_{\phi_f}^{\phi_*} J^2 \frac{\hat{V}_{\text{IG}}}{\hat{V}_{\text{IG},\phi}} \frac{d\phi}{m_{\text{P}}^2}, \quad (3.1)$$

has to be at least enough to resolve the horizon and flatness problems of standard big bang, i.e., [2]

$$\hat{N}_* \simeq 19.4 + 2 \ln \frac{\hat{V}_{\text{IG}}(\phi_*)^{1/4}}{1 \text{ GeV}} - \frac{4}{3} \ln \frac{\hat{V}_{\text{IG}}(\phi_f)^{1/4}}{1 \text{ GeV}} + \frac{1}{3} \ln \frac{T_{\text{rh}}}{1 \text{ GeV}} + \frac{1}{2} \ln \frac{f_{\mathcal{R}}(\phi_*)}{f_{\mathcal{R}}(\phi_f)^{1/3}}, \quad (3.2)$$

where we assumed that IG inflation is followed in turn by a decaying-inflaton, radiation and matter domination, T_{rh} is the reheat temperature after IG inflation, ϕ_* [$\hat{\phi}_*$] is the value of ϕ [$\hat{\phi}$] when k_* crosses outside the inflationary horizon, and ϕ_f [$\hat{\phi}_f$] is the value of ϕ [$\hat{\phi}$] at the end of IG inflation, which can be found, in the slow-roll approximation and for the considered in this paper models, from the condition

$$\max\{\hat{\epsilon}(\phi_f), |\hat{\eta}(\phi_f)|\} = 1, \quad (3.3a)$$

where the slow-roll parameters can be calculated as follows:

$$\hat{\epsilon} = \frac{m_{\text{P}}^2}{2} \left(\frac{\hat{V}_{\text{IG},\hat{\phi}}}{\hat{V}_{\text{IG}}} \right)^2 = \frac{m_{\text{P}}^2}{2J^2} \left(\frac{\hat{V}_{\text{IG},\phi}}{\hat{V}_{\text{IG}}} \right)^2 \text{ and } \hat{\eta} = m_{\text{P}}^2 \frac{\hat{V}_{\text{IG},\hat{\phi}\hat{\phi}}}{\hat{V}_{\text{IG}}} = \frac{m_{\text{P}}^2}{J^2} \left(\frac{\hat{V}_{\text{IG},\phi\phi}}{\hat{V}_{\text{IG}}} - \frac{\hat{V}_{\text{IG},\phi}}{\hat{V}_{\text{IG}}} \frac{J_{,\phi}}{J} \right). \quad (3.3b)$$

3.1.2. The amplitude A_s of the power spectrum of the curvature perturbation generated by ϕ at the pivot scale k_* must to be consistent with data [2]

$$\sqrt{A_s} = \frac{1}{2\sqrt{3}\pi m_{\text{P}}^3} \frac{\hat{V}_{\text{IG}}(\hat{\phi}_*)^{3/2}}{|\hat{V}_{\text{IG},\hat{\phi}}(\hat{\phi}_*)|} = \frac{|J(\phi_*)|}{2\sqrt{3}\pi m_{\text{P}}^3} \frac{\hat{V}_{\text{IG}}(\phi_*)^{3/2}}{|\hat{V}_{\text{IG},\phi}(\phi_*)|} \simeq 4.685 \cdot 10^{-5}, \quad (3.4)$$

where we assume that no other contributions to the observed curvature perturbation exists.

3.1.3. The (scalar) spectral index, n_s , its running, a_s , and the scalar-to-tensor ratio r – estimated through the relations:

$$n_s = 1 - 6\hat{\epsilon}_* + 2\hat{\eta}_*, \quad a_s = 2(4\hat{\eta}_*^2 - (n_s - 1)^2)/3 - 2\hat{\xi}_* \text{ and } r = 16\hat{\epsilon}_*, \quad (3.5)$$

where $\hat{\xi} = m_{\text{P}}^4 \hat{V}_{\text{IG},\hat{\phi}} \hat{V}_{\text{IG},\hat{\phi}\hat{\phi}\hat{\phi}}/\hat{V}^2 = m_{\text{P}}^2 \hat{V}_{\text{IG},\phi} \hat{\eta}_{,\phi}/\hat{V}_{\text{IG}} J^2 + 2\hat{\eta}\hat{\epsilon}$ and the variables with subscript $*$ are evaluated at $\phi = \phi_*$ – must be in agreement with the fitting of the data [2] with ΛCDM model, i.e.,

$$(a) \quad n_s = 0.9603 \pm 0.0146, \quad (b) \quad -0.0314 \leq a_s \leq 0.0046 \quad \text{and} \quad (c) \quad r < 0.11, \quad (3.6)$$

at 95% confidence level (c.l.)

3.1.4. To avoid corrections from quantum gravity and any destabilize of our inflationary scenario due to higher order non-renormalizable terms – see Eq. (2.9) –, we impose two additional theoretical constraints on our models – keeping in mind that $\widehat{V}(\phi_f) \leq \widehat{V}(\phi_*)$:

$$(a) \quad \widehat{V}(\phi_*)^{1/4} \leq m_P \quad \text{and} \quad (b) \quad \phi_* \leq m_P. \quad (3.7)$$

As we show in Sec. 4, the UV cutoff of our model is m_P and so no concerns regarding the validity of the effective theory arise.

3.2 NO-SCALE SUGRA

According to our analysis in Sec. 2.2, IG inflation in the context of no-scale SUGRA can be achieved adopting a Kähler potential which depends at least on two gauge singlet superfields – the inflaton Φ and an accompanying one S – and has the form

$$K = -3m_P^2 \ln \left(\Omega_H(\Phi) + \Omega_H^*(\Phi^*) - \frac{|S|^2}{3m_P^2} + k_S \frac{|S|^4}{3m_P^4} \right), \quad (3.8)$$

as inferred by inserting Eqs. (2.8), (2.7) and (2.5) into Eq. (2.2). Consequently, the Kähler manifold which corresponds to K is $SU(2,1)/SU(2) \times U(1)_R \times \mathbb{Z}_n$ globally symmetric. The underlying symmetry of Kähler manifold allows us to avoid any mixing of inflaton Φ with S which fixes $f_{S\Phi} = 1$ – see Eq. (2.13b). We below extract the inflationary potential in Sec. 3.2.1 and present our analytical and numerical results in Sec. 3.2.2 and 3.2.3 respectively.

3.2.1 THE INFLATIONARY POTENTIAL

Taking into account the form of Ω_H , $f_{\mathcal{R}}$ and $f_{S\Phi}$ from Eqs. (2.9) and (2.13b), Eq. (2.13a) reads

$$\widehat{V}_{IG0} = \frac{\lambda^2 m_P^4 |1 - 2\Omega_H|^2}{4f_{\mathcal{R}}^2} = \frac{\lambda^2 m_P^4 f_{\Phi}^2}{4c_{\mathcal{R}}^4 x_{\phi}^{2n}}, \quad (3.9)$$

since $f_{S\Phi} = 1$ and $f_{\mathcal{R}} = 2c_{\mathcal{R}} x_{\phi}^n / 2^{n/2}$ where we introduce the dimensionless quantities

$$x_{\phi} = \phi / m_P \quad \text{and} \quad f_{\Phi} = 2^{n/2-1} - c_{\mathcal{R}} x_{\phi}^n. \quad (3.10)$$

Obviously \widehat{V}_{IG0} in Eq. (3.9) develops a plateau with almost constant potential energy density corresponding to the Hubble parameter

$$\widehat{H}_{IG} = \frac{\widehat{V}_{IG0}^{1/2}}{\sqrt{3}m_P} \simeq \frac{\lambda m_P}{2\sqrt{3}c_{\mathcal{R}}} \quad \text{with} \quad \widehat{V}_{IG0} \simeq \frac{\lambda^2 m_P^4}{4c_{\mathcal{R}}^2}. \quad (3.11)$$

Along the configuration of Eq. (2.14) $K_{\alpha\bar{\beta}}$ defined in Eq. (2.17a) takes the form

$$(K_{\alpha\bar{\beta}}) = \frac{1}{f_{\mathcal{R}}} \text{diag} \left(\frac{3m_P^2 |\Omega_{H,\phi}|^2}{f_{\mathcal{R}}}, 1 \right) = \text{diag} \left(\frac{3n^2}{2x_{\phi}^2}, \frac{2^{n/2}}{2c_{\mathcal{R}} x_{\phi}^n} \right), \quad (3.12)$$

where the explicit form of Ω_H in Eq. (2.9) is taken into account. Integrating the first equation in Eq. (3.12) we can identify the EF field:

$$\widehat{\phi} = \widehat{\phi}_c + \sqrt{\frac{3}{2}} n m_P \ln \frac{\phi}{\langle \phi \rangle} \quad \text{with} \quad \langle \phi \rangle = \frac{\sqrt{2} m_P}{\sqrt[n]{2c_{\mathcal{R}}}}, \quad (3.13)$$

where we take into account Eqs. (2.9) and (2.12b). Also $\widehat{\phi}_c$ is a constant of integration.

Following the general analysis in Sec. 2.3 we derive the mass spectrum along the configuration of Eq. (2.14). Our results are arranged in Table 1. We see there that $k_S \gtrsim 1$ assists us to achieve $\widehat{m}_s^2 > 0$ – in accordance with Ref. [15, 17, 18]. Inserting the extracted masses in Eq. (2.18) we can proceed to the numerical analysis of IG inflation in the EF [4], employing the standard slow-roll approximation [25] – see Sec. 3.2.3. For the sake of the presentation, however, we first – see Sec. 3.2.2 – present analytic results based on Eq. (3.11), which are quite close to the numerical ones.

FIELDS	EINGESTATES	MASSES SQUARED
1 real scalar	$\hat{\theta}$	$\hat{m}_{\theta}^2 = \lambda^2 m_{\text{P}}^2 (2^{n-2} - c_{\mathcal{R}}^2 x_{\phi}^n f_{\Phi}) / 3c_{\mathcal{R}}^4 x_{\phi}^{2n} \simeq 4\hat{H}_{\text{IG}}^2$
2 real scalars	$\hat{s}, \hat{\bar{s}}$	$\hat{m}_s^2 = \lambda^2 m_{\text{P}}^2 (2^{3n/2} + 4c_{\mathcal{R}} x_{\phi}^n (2^n - 2^{n/2} c_{\mathcal{R}} x_{\phi}^n + 12k_S f_{\Phi}^2)) / 3 \cdot 2^{3+n/2} c_{\mathcal{R}}^4 x_{\phi}^{2n}$
2 Weyl spinors	$\hat{\psi}_{\pm} = \frac{\hat{\psi}_{\Phi} \pm \hat{\psi}_S}{\sqrt{2}}$	$\hat{m}_{\psi\pm}^2 \simeq 2^{n-2} \lambda^2 m_{\text{P}}^2 / 3c_{\mathcal{R}}^4 x_{\phi}^{2n}$

Table 1: The mass spectrum along the trajectory of Eq. (2.14) during IG inflation.

3.2.2 ANALYTICAL RESULTS

The duration of the slow-roll IG inflation is controlled by the slow-roll parameters which, according to their definition in Eq. (3.3b), are calculated to be

$$\hat{\epsilon} \simeq \frac{2^n}{3f_{\Phi}^2} \quad \text{and} \quad \hat{\eta} \simeq \frac{2^{1+n/2}(2^{n/2} - c_{\mathcal{R}} x_{\phi}^n)}{3f_{\Phi}^2}. \quad (3.14)$$

The termination of IG inflation is triggered by the violation of the $\hat{\epsilon}$ criterion at $\phi = \phi_f$ given by

$$\hat{\epsilon}(\phi_f) = 1 \Rightarrow \phi_f = \sqrt{2}m_{\text{P}} \left((\sqrt{3} + 2)/2\sqrt{3}c_{\mathcal{R}} \right)^{1/n}, \quad (3.15a)$$

since the violation of the $\hat{\eta}$ criterion occurs at $\phi = \tilde{\phi}_f$ such that

$$\hat{\eta}(\tilde{\phi}_f) = 1 \Rightarrow \tilde{\phi}_f = \sqrt{2}m_{\text{P}} \left(\frac{5}{6c_{\mathcal{R}}} \right)^{1/n} = \left((3 + 2\sqrt{3})/5 \right)^{-1/n} \phi_f < \phi_f. \quad (3.15b)$$

In the EF, $\hat{\phi}_f$ remains independent of $c_{\mathcal{R}}$ and n , since substituting Eq. (3.15a) into Eq. (3.13) we obtain

$$\hat{\phi}_f - \hat{\phi}_c \simeq \sqrt{3/2}m_{\text{P}} \ln(1 + 2/\sqrt{3}). \quad (3.16)$$

E.g., setting $\hat{\phi}_c = 0$, we obtain $\hat{\phi}_f = 0.94m_{\text{P}}$.

Given that $\phi_f \ll \phi_*$, we can find a relation between ϕ_* and \hat{N}_* as follows

$$\hat{N}_* \simeq \frac{3c_{\mathcal{R}}}{2^{1+n/2}m_{\text{P}}^n} (\phi_*^n - \phi_f^n) \Rightarrow \phi_* \simeq m_{\text{P}} \sqrt[n]{2^{1+n/2}\hat{N}_*/3c_{\mathcal{R}}}. \quad (3.17a)$$

Obviously, IG inflation consistent with Eq. (3.7b) can be achieved if

$$\phi_* \leq m_{\text{P}} \Rightarrow c_{\mathcal{R}} \geq 2^{1+n/2}\hat{N}_*/3 \quad (3.17b)$$

for $\hat{N}_* \simeq 52$. Therefore, we need relatively large $c_{\mathcal{R}}$'s which increase with n . On the other hand, $\hat{\phi}_*$ remains transplanckian, since plugging Eq. (3.17a) into Eq. (3.13) we find

$$\hat{\phi}_* \simeq \hat{\phi}_c + \sqrt{3/2}m_{\text{P}} \ln(4\hat{N}_*/3), \quad (3.18)$$

which gives $\hat{\phi}_* = 5.3m_{\text{P}}$ for $\hat{\phi}_c = 0$. Despite this fact, our construction remains stable under possible corrections from non-renormalizable terms in Ω_{H} since these are expressed in terms of initial field Φ , and can be harmless for $|\Phi| \leq m_{\text{P}}$.

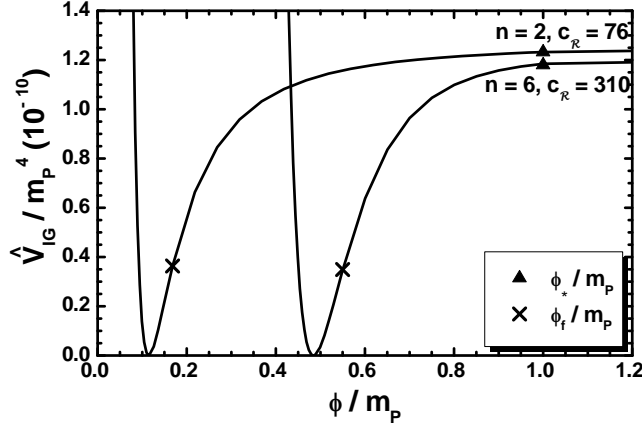


Figure 1: The inflationary potential \hat{V}_{IG} as a function of ϕ for $n = 2, \lambda = 1.7 \cdot 10^{-3}$ and $c_{\mathcal{R}} = 76$ or $n = 6, \lambda = 6.8 \cdot 10^{-3}$ and $c_{\mathcal{R}} = 310$. The values corresponding to ϕ_* and ϕ_f are also depicted.

Upon substitution of Eqs. (3.11), (3.12) and (3.17a) into Eq. (3.4) we find A_s as follows

$$A_s^{1/2} = \frac{\lambda m_{\text{P}}^n f_{\Phi}(\phi_*)^2}{2^{n/2+2} \sqrt{2} \pi c_{\mathcal{R}}^2 \phi_*^n} = \frac{\lambda(3 - 4\hat{N}_*)^2}{96 \sqrt{2} c_{\mathcal{R}} \pi \hat{N}_*} \Rightarrow \lambda \simeq 6\pi \sqrt{2 A_s c_{\mathcal{R}} / \hat{N}_*} \Rightarrow c_{\mathcal{R}} \simeq 41637\lambda, \quad (3.19)$$

for $\hat{N}_* \simeq 52$. Therefore, enforcing Eq. (3.4) we obtain a relation between λ and $c_{\mathcal{R}}$ which turns out to be independent of n . Replacing ϕ_* by Eq. (3.17a) into Eq. (3.5) we estimate, finally, the inflationary observable through the relations:

$$n_s = \frac{(1 + 4\hat{N}_*)(4\hat{N}_* - 15)}{(3 - 4\hat{N}_*)^2} \simeq 1 - 2/\hat{N}_* - 9/2\hat{N}_*^2 = 0.960, \quad (3.20a)$$

$$a_s \simeq -2\hat{\xi}_* = \frac{128(3 - \hat{N}_*)}{(4\hat{N}_* - 3)^3} \simeq -2/\hat{N}_*^2 + 3/2\hat{N}_*^3 = -0.0007, \quad (3.20b)$$

$$r = \frac{192}{(3 - 4\hat{N}_*)^2} \simeq 12/\hat{N}_*^2 = 0.0045 \quad (3.20c)$$

for $\hat{N}_* \simeq 52$. These outputs are fully consistent with the observational data, Eq. (3.6).

3.2.3 NUMERICAL RESULTS

The inflationary scenario under consideration depends on the parameters:

$$\lambda, c_{\mathcal{R}}, k_S \text{ and } T_{\text{rh}}.$$

Our results are essentially independent of k_S 's, provided that we choose them so as $\hat{m}_s^2 > 0$ for every allowed λ and $c_{\mathcal{R}}$ – see Table 1. We therefore set $k_S = 1$ throughout our calculation. We also choose $\Lambda \simeq 10^{13}$ GeV so as the one-loop corrections in Eq. (2.18) vanish at the SUSY vacuum, Eqs. (2.12b) and (2.6). Finally we choose $T_{\text{rh}} = 10^9$ GeV as suggested by reliable post-inflationary scenario – see Ref. [11]. Upon substitution of \hat{V}_{IG} from Eqs. (2.18) and (3.11) in Eqs. (3.3b), (3.1) and (3.4) we extract the inflationary observables as functions of $c_{\mathcal{R}}$, λ and ϕ_* . The two latter parameters can be determined by enforcing the fulfilment of Eq. (3.2) and (3.4), for every chosen $c_{\mathcal{R}}$. Our numerical findings are quite close to the analytic ones listed in Sec. 3.2.2 for presentational purposes.

The variation of \hat{V}_{IG} as a function of ϕ for two different values of n can be easily inferred from Fig. 1, where we depict \hat{V}_{IG} versus ϕ for $\phi_* = m_{\text{P}}$ and $n = 2$ or $n = 6$. The imposition $\phi_* = m_{\text{P}}$

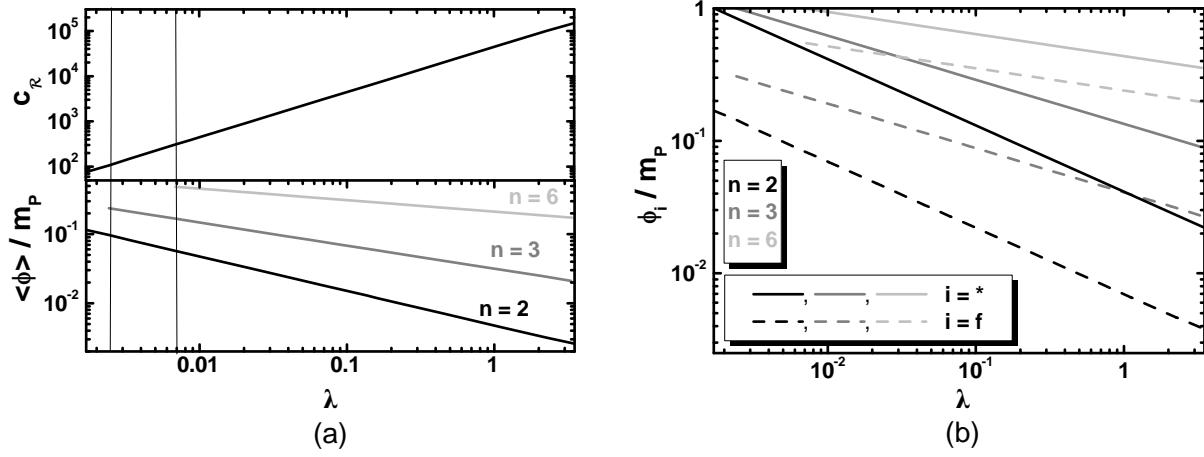


Figure 2: The allowed by Eqs. (3.2), (3.4) and (3.7) values of $c_{\mathcal{R}}$ and the resulting $\langle \phi \rangle$ [ϕ_* (solid line) and ϕ_f (dashed line)] versus λ (a) [(b)]. We use black, gray and light gray lines for $n = 2, 3$ and 6 respectively, $k_S = 1$ and $T_{\text{rh}} = 10^9$ GeV. Eq. (3.7) is fulfilled to the right of the thin line.

correspond to $\lambda = 0.0017$ and $c_{\mathcal{R}} = 76$ for $n = 2$ and $\lambda = 0.0068$ and $c_{\mathcal{R}} = 310$ for $n = 6$. In accordance with our findings in Eqs. (3.13) and (3.17b) we conclude that increasing n (i) larger $c_{\mathcal{R}}$'s and therefore lower \hat{V}_{IG0} 's are required to obtain $\phi < m_P$; (ii) larger ϕ_f and $\langle \phi \rangle$ are obtained. Combining Eqs. (3.15a) and (3.19) with Eq. (3.11) we can convince ourselves that $\hat{V}_{\text{IG0}}(\phi_f)$ is independent of $c_{\mathcal{R}}$ and to a considerable degree of n .

By varying λ we can delineate the region of the parameters allowed by a simultaneous imposition of Eqs. (3.4), (3.2) and (3.7). Our results are displayed in Fig. 2, where we draw as functions of λ the allowed values of $c_{\mathcal{R}}$ and $\langle \phi \rangle$ – see Fig. 2-(a) – ϕ_* (solid line) and ϕ_f (dashed line) – see Fig. 2-(b). We use black, gray and light gray lines for $n = 2, 3$ and 6 respectively. As anticipated in Eq. (3.19) the relation between $c_{\mathcal{R}}$ and λ is independent of n ; the various lines, thus, coincide. However, Eq. (3.7) is fulfilled to the right of the thin line. Indeed, the lower bound of the depicted lines comes from the saturation of Eq. (3.17b) whereas the upper bound originates from the perturbative bound on λ , $\lambda \leq \sqrt{4\pi} \simeq 3.54$. Moreover, the variation of ϕ_f and ϕ_* as a function of λ – drawn in Fig. 2-(b) – is consistent with Eqs. (3.15a) and (3.17a).

The overall allowed parameter space of the model for $n = 2, 3$ and 6 is correspondingly

$$76, 105, 310 \lesssim c_{\mathcal{R}} \lesssim 1.5 \cdot 10^5 \quad \text{and} \quad (1.7, 2.4, 6.8) \cdot 10^{-3} \lesssim \lambda \lesssim 3.54 \quad \text{for} \quad \hat{N}_* \simeq 52 \quad (3.21a)$$

with $\langle \phi \rangle$ being confined in the ranges $(0.0026 - 0.1)$, $(0.021 - 0.24)$ and $(0.17 - 0.48)$. Moreover, the masses of the various scalars in Table 1 remain well above \hat{H}_{IG} both during and after IG inflation for the selected k_S . E.g., for $n = 3$ and $c_{\mathcal{R}} = 495$ (corresponding to $\lambda = 0.01$) we obtain

$$(\hat{m}_{\theta}^2(\phi_*), \hat{m}_s^2(\phi_*)) / \hat{H}_{\text{IG}}^2(\phi_*) \simeq (4, 905) \quad \text{and} \quad (\hat{m}_{\theta}^2(\phi_f), \hat{m}_s^2(\phi_f)) / \hat{H}_{\text{IG}}^2(\phi_f) \simeq (10.5, 26.8). \quad (3.21b)$$

Letting λ or $c_{\mathcal{R}}$ vary within its allowed region in Eq. (3.21a), independently of n , we obtain

$$0.961 \lesssim n_s \lesssim 0.963, \quad -7 \lesssim a_s / 10^{-4} \lesssim -6.4 \quad \text{and} \quad 4.2 \gtrsim r / 10^{-3} \gtrsim 3.6, \quad (3.22)$$

which lie close to the analytic results in Eqs. (3.20a), (3.20b) and (3.20c) and within the allowed ranges of Eq. (3.6), with n_s being impressively spot on its central observationally favored value – see Eq. (3.6a). Therefore, the inclusion of the variant exponent $n \geq 2$, compared to the initial model of Ref. [11], does not affect the successful predictions on the inflationary observables.

3.3 BEYOND NO-SCALE SUGRA

If we lift the assumption of no-scale SUGRA in Eq. (2.8), Ω takes its more general form, obtained by inserting Eqs. (2.7) and (2.9) into Eq. (2.5); the resulting through Eq. (2.2) Kähler potential is

$$K = -3m_{\text{P}}^2 \ln \left(\Omega_{\text{H}}(\Phi) + \Omega_{\text{H}}^*(\Phi^*) - \frac{|S|^2}{3m_{\text{P}}^2} - \frac{|\Phi|^2}{3m_{\text{P}}^2} + k_S \frac{|S|^4}{3m_{\text{P}}^4} + 2k_{\Phi} \frac{|\Phi|^4}{3m_{\text{P}}^4} + 2k_{S\Phi} \frac{|S|^2|\Phi|^2}{3m_{\text{P}}^4} \right), \quad (3.23)$$

where the factors of 2 are added just for convenience. The description of the inflationary potential, our analytical and numerical results are exhibited below in Secs. 3.3.1, 3.3.2 and 3.3.3 correspondingly.

3.3.1 THE INFLATIONARY POTENTIAL

The tree-level scalar potential in this case has its general form in Eq. (2.13a) where $f_{\mathcal{R}}$ and $f_{S\Phi}$ are calculated by employing their definitions in Eq. (2.13b) as follows

$$f_{\mathcal{R}} = 2c_{\mathcal{R}} \frac{x_{\phi}^n}{2^{n/2}} + \frac{x_{\phi}^2}{6} + \frac{k_{\Phi}}{12} x_{\phi}^4 \quad \text{and} \quad f_{S\Phi} = 1 - k_{S\Phi} x_{\phi}^2. \quad (3.24)$$

Taking into account the form of $f_{\mathcal{R}}$ above, \widehat{V}_{IG0} can be cast as follows

$$\widehat{V}_{\text{IG0}} = \frac{\lambda^2 f_{\Phi}^2}{4c_{\mathcal{R}}^2 f_{S\Phi} x_{\phi}^4 (c_{\mathcal{R}} x_{\phi}^{n-2} - 2^{n/2-2} f_{\phi\phi}/3)^2} \quad (3.25a)$$

where $f_{\phi\phi} = 1 - k_{\Phi} x_{\phi}^2$ while x_{ϕ} and f_{Φ} are defined in Eq. (3.10). Similarly to Sec. 3.2, \widehat{V}_{IG0} in Eq. (3.25a) develops a plateau with almost constant potential energy density corresponding to the Hubble parameter

$$\widehat{H}_{\text{IG}} = \frac{\widehat{V}_{\text{IG0}}^{1/2}}{\sqrt{3}m_{\text{P}}} \simeq \frac{\lambda m_{\text{P}}}{2\sqrt{3}f_{S\Phi}c_{\mathcal{R}}} \quad \text{with} \quad \widehat{V}_{\text{IG0}} \simeq \frac{\lambda^2 m_{\text{P}}^4}{4f_{S\Phi}c_{\mathcal{R}}^2}. \quad (3.25b)$$

Moreover, the EF canonically normalized inflaton, $\widehat{\phi}$, is found via Eq. (2.17b) with J^2 given by

$$J^2 = \frac{3n^2 c_{\mathcal{R}}^2 x_{\phi}^{2n} + 2^{4+n/2} c_{\mathcal{R}} x_{\phi}^{2+n} (1 - n + 2k_{\Phi}(n-2)x_{\phi}^2)}{(c_{\mathcal{R}} x_{\phi}^{1+n} - 2^{n/2-2} x_{\phi}^3 f_{\phi\phi}/3)^2} \simeq \frac{3n^2}{2x_{\phi}^2} + \frac{2^{n/2}(1-n)}{2c_{\mathcal{R}} x_{\phi}^n}. \quad (3.26)$$

Consequently, J turns out to be close to that obtained in Sec. 3.2.1.

Following the standard procedure of Sec. 2.3 we construct the mass spectrum of the theory along the path of Eq. (2.14). The precise expressions of the relevant masses squared, taken into account in our numerical computation, are rather lengthy due to the numerous contributions to \widehat{V}_{IG0} , Eq. (3.25a). Our findings, though, can be considerably simplified, if we perform an expansion for small x_{ϕ} 's – retaining f_{Φ} intact –, consistently with our restriction, Eq. (3.7). If we keep the lowest order terms, the masses squared for the scalars reduce to those displayed in Table 1, whereas the mass squared of the chiral fermions shown in Table 1 has to be multiplied by the factor

$$1 + k_{S\Phi} c_{\mathcal{R}} x_{\phi}^{2+n} / 2^{n/2-1} n. \quad (3.27)$$

As in the case of Sec. 3.2, employing the mass spectrum along the direction of Eq. (2.14), we can calculate \widehat{V}_{IG} in Eq. (2.18) to further analyze the model.

3.3.2 ANALYTICAL RESULTS

Upon substitution of Eq. (3.25b) into Eq. (3.3b), we can extract the slow-roll parameters which determine the strength of the inflationary stage. Performing expansions about $x_\phi \simeq 0$, as above, we can extract approximate expressions which assist us to interpret the numerical results presented in Sec. 3.3.3. Namely, we find

$$\hat{\epsilon} = \frac{(2^{n/2}n + 2k_{S\Phi}c_{\mathcal{R}}x_\phi^{2+n})^2}{3n^2f_\Phi^2}, \quad (3.28a)$$

$$\hat{\eta} = \left(2^n n^2 + 4k_{S\Phi}c_{\mathcal{R}}^2 x_\phi^{2(1+n)} + 2^{n/2}c_{\mathcal{R}}x_\phi^n \left(((n-2)^2/6 + 4k_{S\Phi}(n-1))x_\phi^2 - n^2 \right) \right) / 3n^2 f_\Phi^2. \quad (3.28b)$$

As it may be numerically verified, ϕ_* and ϕ_f do not decline a lot from their values in Eqs. (3.17a) and (3.15a), which can be served for our estimations below. In particular, replacing \hat{V}_{IG0} from Eq. (3.25b) in Eq. (3.4) we obtain

$$A_s^{1/2} = \frac{n\lambda m_{\text{P}}^n f_\Phi^2(\phi_*)}{4\sqrt{2}\pi c_{\mathcal{R}}^2 \phi_*^n (2^{n/2}n + 2k_{S\Phi}c_{\mathcal{R}}x_\phi^{2+n})} \Rightarrow \lambda \simeq 2\pi\sqrt{2A_s}c_{\mathcal{R}} \left(\frac{3}{\hat{N}_*} + \frac{8k_{S\Phi}}{n} \left(\frac{2\hat{N}_*}{3c_{\mathcal{R}}} \right)^{2/n} \right). \quad (3.29)$$

Comparing this expression with the one obtained in the case of no-scale SUGRA, Eq. (3.19), we remark that λ acquires a mild dependence on both $k_{S\Phi}$ and n . Inserting, Eq. (3.17a) into Eq. (3.5) we can similarly provide an expression for n_s . This is

$$n_s \simeq 1 - 2/\hat{N}_* + \left(\frac{4}{9} \right)^{1/n} \left(\frac{\hat{N}_*}{c_{\mathcal{R}}} \right)^{2/n} \frac{128k_{S\Phi} + 27n^2/\hat{N}_*^3}{12n^2}. \quad (3.30)$$

Therefore, a clear dependence of n_s on n and $k_{S\Phi}$ arises, with the second one being much more efficient. On the other hand, a_s and r remain pretty close to those obtained in the absence of $k_{S\Phi}$ – see Sec. 3.2.2.

3.3.3 NUMERICAL RESULTS

This inflationary scenario depends on the following parameters:

$$\lambda, c_{\mathcal{R}}, k_S, k_{S\Phi}, k_\Phi \text{ and } T_{\text{rh}}.$$

As in the case of Sec. 3.2.3 our results are independent of k_S , provided that $\hat{m}_s^2 > 0$ – see in Table 1. The same is also valid for k_Φ since the contribution from the second term in $f_{\mathcal{R}}$, Eq. (3.24), is overshadowed by the strong enough first term including $c_{\mathcal{R}} \gg 1$. We therefore set $k_S = 1$ and $k_\Phi = 0.5$. We also choose $T_{\text{rh}} = 10^9$ GeV. Besides these values, in our numerical code, we use as input parameters $c_{\mathcal{R}}$, $k_{S\Phi}$ and ϕ_* . For every chosen $c_{\mathcal{R}} \geq 1$, we restrict λ and ϕ_* so that the conditions Eqs. (3.1), (3.4) and (3.7) are satisfied. By adjusting $k_{S\Phi}$ we can achieve n_s 's in the range of Eq. (3.6). Our results are displayed in Fig. 3-(a₁) and (a₂) [Fig. 3-(b₁) and (b₂)], where we delineate the hatched regions allowed by Eqs. (3.1), (3.4), (3.6) and (3.7) in the $\lambda - c_{\mathcal{R}}$ [$\lambda - k_{S\Phi}$] plane. We take $n = 2$ in Fig. 3-(a₁) and (b₁) and $n = 3$ in Fig. 3-(a₂) and (b₂). The conventions adopted for the various lines are also shown. In particular, the dashed [dot-dashed] lines correspond to $n_s = 0.975$ [$n_s = 0.946$], whereas the solid (thick) lines are obtained by fixing $n_s = 0.96$ – see Eq. (3.6). Along the thin line, which provides the lower bound for the regions presented in Fig. 3, the constraint of Eq. (3.7b) is saturated. At the other end, the perturbative bound on λ bounds the various regions.

From Fig. 3-(a₁) and (a₂) we see that $c_{\mathcal{R}}$ remains almost proportional to λ and for constant λ , $c_{\mathcal{R}}$ increases as n_s decreases. From Fig. 3-(b₁) we remark that $k_{S\Phi}$ is confined close to zero for $n_s = 0.96$

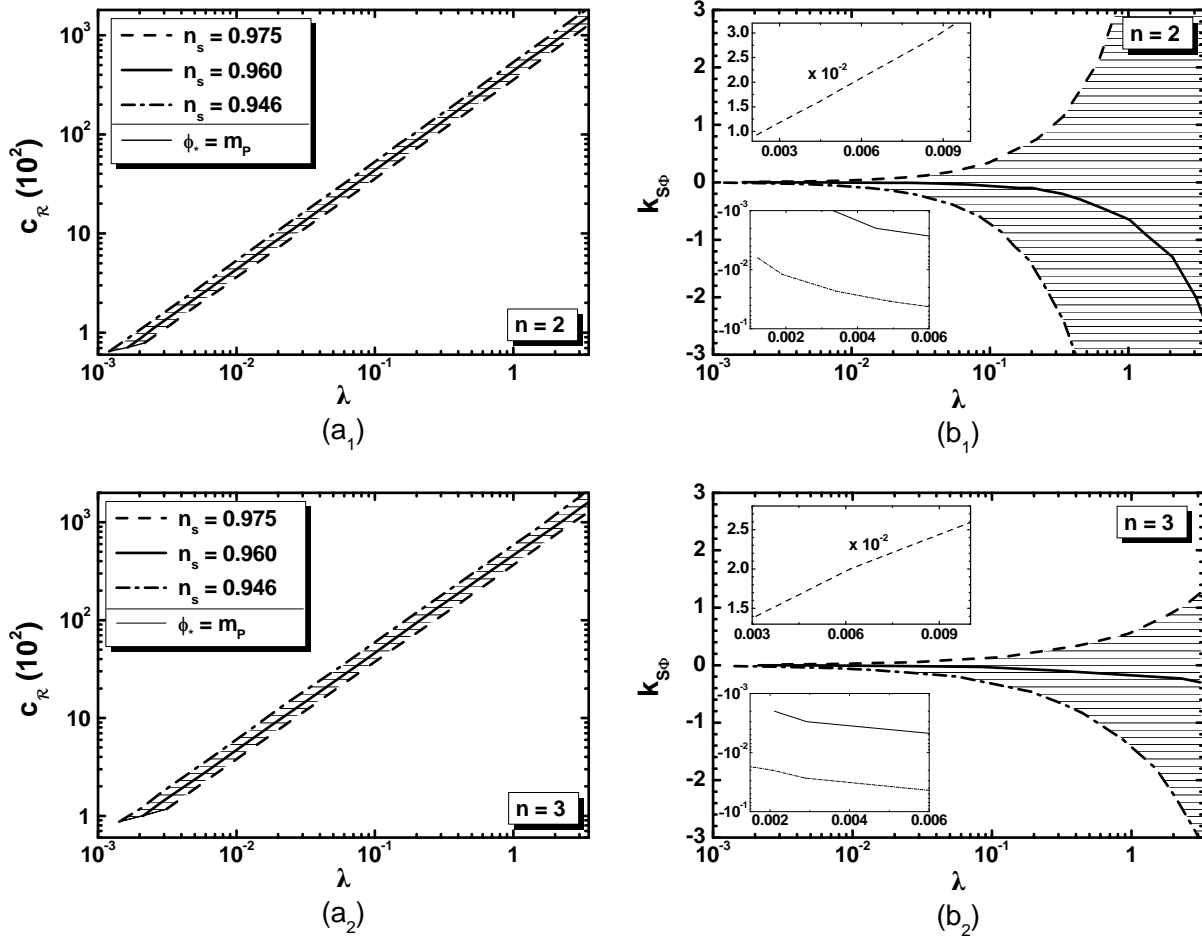


Figure 3: The (hatched) regions allowed by Eqs. (3.2), (3.4), (3.6) and (3.7) in the $\lambda - c_{\mathcal{R}}$ plane (a_1 , a_2) and $\lambda - k_{S\Phi}$ plane (b_1 , b_2) for $k_S = 1$, $k_{\Phi} = 0.5$ and $n = 2$ (a_1 , b_1) or $n = 3$ (a_2 , b_2). The conventions adopted for the various lines are also shown.

and $\lambda < 0.16$ or $\phi_* > 0.1 m_P$ – see Eq. (3.17a). Therefore, a degree of tuning (of the order of 10^{-2}) is needed in order to reproduce the experimental data of Eq. (3.6a). On the other hand, for $\lambda > 0.16$ (or $\phi_* < 0.1 m_P$), $k_{S\Phi}$ takes quite natural (of order one) negative values – consistently with Eq. (3.30). This feature, however, does not insist for $n = 3$ – see Fig. 3-(b_2) –, where the allowed (hatched) region is considerably shrunk and so, $k_{S\Phi}$ remains constantly below unity for any λ . As we explicitly verified, for $n = 6$ the results turn out to be even more concentrated about $k_{S\Phi} \simeq 0$. Therefore, we can conclude that this embedding of IG inflation in SUGRA favors low n values.

More explicitly, for $n_s = 0.96$ and $\hat{N}_* \simeq 52$ we find:

$$71 \lesssim c_{\mathcal{R}} \lesssim 1.5 \cdot 10^5 \quad \text{with} \quad 1.6 \cdot 10^{-3} \lesssim \lambda \lesssim 3.5 \quad \text{and} \quad 0 \lesssim -k_{S\Phi} \lesssim 2.4 \quad (n = 2); \quad (3.31a)$$

$$100 \lesssim c_{\mathcal{R}} \lesssim 1.4 \cdot 10^5 \quad \text{with} \quad 2.1 \cdot 10^{-3} \lesssim \lambda \lesssim 3.5 \quad \text{and} \quad 0.002 \lesssim -k_{S\Phi} \lesssim 0.3 \quad (n = 3); \quad (3.31b)$$

$$270 \lesssim c_{\mathcal{R}} \lesssim 1.65 \cdot 10^5 \quad \text{with} \quad 5.6 \cdot 10^{-3} \lesssim \lambda \lesssim 3.5 \quad \text{and} \quad 0.01 \lesssim -k_{S\Phi} \lesssim 0.1 \quad (n = 6). \quad (3.31c)$$

Note that the lower bounds on $c_{\mathcal{R}}$ and λ are quite close to those obtained in Eq. (3.21a). In both cases $6.8 \lesssim |a_s|/10^{-4} \lesssim 8.2$ and $r \simeq 3.8 \cdot 10^{-3}$ which lie within the allowed ranges of Eq. (3.6). Needless to say that, as in Sec. 3.2.3, we here also obtain $\hat{m}_{\chi\alpha}^2/\hat{H}_{\text{IG}}^2 \gg 1$ with $\hat{m}_{\chi\alpha}^2$ being defined in Eq. (2.16a).

4 THE EFFECTIVE CUT-OFF SCALE

An outstanding trademark of IG inflation is that it is unitarity-safe, despite the fact that its implementation with subplanckian ϕ 's – see Eq. (3.17b) – requires relatively large $c_{\mathcal{R}}$'s. To show this we below extract the UV cut-off scale, Λ_{UV} , of the effective theory first in the JF – Sec. 4.1 – and then in the EF – see Sec. 4.2. Although the expansions about $\langle\phi\rangle$ presented below are not valid [9] during IG inflation, we consider the extracted this way Λ_{UV} as the overall cut-off scale of the theory, since reheating is an unavoidable stage of the inflationary dynamics [10].

4.1 JORDAN FRAME COMPUTATION

The possible problematic process in the JF, which causes [8] concerns about the unitarity-violation, is the $\delta\phi - \delta\phi$ scattering process via s -channel graviton, $h^{\mu\nu}$, exchange – $\bar{\delta\phi}$ represents an excitation of ϕ about $\langle\phi\rangle$, see below. The relevant vertex is $c_{\mathcal{R}}\delta\phi^2\Box h/m_{\text{P}}$ – with $h = h_{\mu}^{\mu}$ – can be derived from the first term in the right-hand side of Eq. (2.3) expanding the JF metric $g_{\mu\nu}$ about the flat spacetime metric $\eta_{\mu\nu}$ and the inflaton ϕ about its v.e.v as follows:

$$g_{\mu\nu} \simeq \eta_{\mu\nu} + h_{\mu\nu}/m_{\text{P}} \quad \text{and} \quad \phi = \langle\phi\rangle + \delta\phi. \quad (4.1)$$

Retaining only the terms with two derivatives of the excitations, the part of the lagrangian corresponding to the two first terms in the right-hand side of Eq. (2.3) takes the form

$$\begin{aligned} \delta\mathcal{L} &= -\frac{\langle\Omega_{\text{H}}\rangle}{4}F_{\text{EH}}(h^{\mu\nu}) + \frac{1}{2}\partial_{\mu}\delta\phi\partial^{\mu}\delta\phi + \left(m_{\text{P}}\langle\Omega_{\text{H},\phi}\rangle + \delta_{\mathcal{R}}c_{\mathcal{R}}^{2/n}\frac{\delta\phi}{m_{\text{P}}}\right)F_{\mathcal{R}}\delta\phi + \dots \\ &= -\frac{1}{8}F_{\text{EH}}(\bar{h}^{\mu\nu}) + \frac{1}{2}\partial_{\mu}\bar{\delta\phi}\partial^{\mu}\bar{\delta\phi} + \delta_{\mathcal{R}}\frac{c_{\mathcal{R}}^{2/n}}{\sqrt{2}m_{\text{P}}}\frac{\sqrt{\langle\Omega_{\text{H}}\rangle}}{\langle\Omega_{\text{H}}\rangle}\bar{\delta\phi}^2\Box\bar{h} + \dots, \end{aligned} \quad (4.2a)$$

where $\delta_{\mathcal{R}} = 1/2$ [$\delta_{\mathcal{R}} = 2^{2/n}n(n-1)/8$] for $n = 2$ [$n > 2$] and the functions F_{EH} and $F_{\mathcal{R}}$ read

$$F_{\text{EH}}(h^{\mu\nu}) = h^{\mu\nu}\Box h_{\mu\nu} - h\Box h + 2\partial_{\rho}h^{\mu\rho}\partial^{\nu}h_{\mu\nu} - 2\partial_{\nu}h^{\mu\nu}\partial_{\mu}h, \quad (4.2b)$$

$$F_{\mathcal{R}}(h^{\mu\nu}) = \Box h - \partial_{\mu}\partial_{\nu}h^{\mu\nu}. \quad (4.2c)$$

The JF canonically normalized fields $\bar{h}_{\mu\nu}$ and $\bar{\delta\phi}$ are defined by the relations

$$\bar{\delta\phi} = \sqrt{\frac{\langle\Omega_{\text{H}}\rangle}{\langle\Omega_{\text{H}}\rangle}}\delta\phi \quad \text{and} \quad \frac{\bar{h}_{\mu\nu}}{\sqrt{2}} = \sqrt{\langle\Omega_{\text{H}}\rangle}h_{\mu\nu} + \frac{m_{\text{P}}\langle\Omega_{\text{H},\phi}\rangle}{\sqrt{\langle\Omega_{\text{H}}\rangle}}\eta_{\mu\nu}\delta\phi \quad \text{with} \quad \bar{\Omega}_{\text{H}} = \Omega_{\text{H}} + 3m_{\text{P}}^2\Omega_{\text{H},\phi}^2. \quad (4.2d)$$

The interaction originating from the last term in the right-hand side of Eq. (4.2a) gives rise to a scattering amplitude which is written in terms of the center-of-mass energy E as follows

$$\mathcal{A} \sim \left(\frac{E}{\Lambda_{\text{UV}}}\right)^2 \quad \text{with} \quad \Lambda_{\text{UV}} = \frac{m_{\text{P}}}{\delta_{\mathcal{R}}c_{\mathcal{R}}^{2/n}}\frac{\langle\bar{\Omega}_{\text{H}}\rangle}{\sqrt{\langle\Omega_{\text{H}}\rangle}} = \frac{m_{\text{P}}}{\delta_{\mathcal{R}}c_{\mathcal{R}}^{2/n}}\left(\sqrt{\langle\Omega_{\text{H}}\rangle} + 3m_{\text{P}}^2\frac{\langle\Omega_{\text{H},\phi}\rangle^2}{\sqrt{\langle\Omega_{\text{H}}\rangle}}\right) \simeq m_{\text{P}} \quad (4.3)$$

(up to irrelevant numerical prefactors) since $\langle\Omega_{\text{H}}\rangle = 1/2 \ll m_{\text{P}}^2\langle\Omega_{\text{H},\phi}\rangle^2 \simeq 2^{2/n}n^2c_{\mathcal{R}}^{2/n}/8$ and Λ_{UV} is identified as the UV cut-off scale in the JF, since \mathcal{A} remains within the validity of the perturbation theory provided that $E < \Lambda_{\text{UV}}$. Obviously, the argument above can be equally well applied to both implementations of IG inflation in SUGRA – see Sec. 3.2 and 3.3 – since the extra terms included in Eq. (3.23) – compared to Eq. (3.8) – are small enough and do not generate any problem with the perturbative unitarity.

4.2 EINSTEIN FRAME COMPUTATION

Alternatively, Λ_{UV} can be determined in EF, following the systematic approach of Ref. [10]. Note, in passing, that the EF (canonically normalized) inflaton,

$$\widehat{\delta\phi} = \langle J \rangle \delta\phi \quad \text{with} \quad \langle J \rangle = \sqrt{\frac{3}{2}} \frac{n}{\langle x_\phi \rangle} = \frac{\sqrt{3}}{2} n \sqrt[2]{c_{\mathcal{R}}} \quad (4.4)$$

acquires mass which is given by

$$\widehat{m}_{\delta\phi} = \left\langle \widehat{V}_{IG0, \widehat{\phi\phi}} \right\rangle^{1/2} = \left\langle \widehat{V}_{IG0, \phi\phi} / J^2 \right\rangle^{1/2} = \lambda m_P / \sqrt{3} c_{\mathcal{R}}. \quad (4.5)$$

Making use of Eq. (3.19) we find $\widehat{m}_{\delta\phi} = 3 \cdot 10^{13}$ GeV for the case of no-scale SUGRA independently of the value of n – in accordance with the findings in Ref. [12]. Beyond no-scale SUGRA, replacing λ in Eq. (4.5) from Eq. (3.29), we find that $\widehat{m}_{\delta\phi}$ inherits from λ a mild dependence on both n and $k_{S\Phi}$. E.g., for $\phi_* = 0.5 m_P$, $n = 2 - 6$ and n_s in the range of Eq. (3.6) we find $2.2 \lesssim \widehat{m}_{\delta\phi} / 10^{13} \text{ GeV} \lesssim 3.8$ with the lower [upper] value corresponding to the lower [upper] bound on n_s in Eq. (3.6) – see Fig. 3-(a₁) and (a₂).

The fact that $\widehat{\delta\phi}$ does not coincide with $\delta\phi$ – contrary to the standard Higgs inflation [8,9] – ensures that the IG models are valid up to m_P . To show it, we write the EF action S in Eq. (2.1a) along the path of Eq. (2.14) as follows

$$S = \int d^4x \sqrt{-\widehat{g}} \left(-\frac{1}{2} m_P^2 \widehat{\mathcal{R}} + \frac{1}{2} J^2 \dot{\phi}^2 - \widehat{V}_{IG0} + \dots \right), \quad (4.6a)$$

where the dot denotes derivation w.r.t the JF cosmic time and the ellipsis represents terms irrelevant for our analysis. Also J and \widehat{V}_{IG0} are respectively given by Eqs. (2.17b) and (3.11) [Eqs. (3.26) and (3.25b)] for the model of Sec. 3.2 [Sec. 3.3]. For both models, J^2 is accurately enough estimated by Eq. (3.12) – cf. Eq. (3.26). Expanding $J^2 \dot{\phi}^2$ about $\langle \phi \rangle$ – see Eq. (3.13) – in terms of $\widehat{\delta\phi}$ in Eq. (4.4) we arrive at the following result

$$J^2 \dot{\phi}^2 = \left(1 - \frac{2}{n} \sqrt{\frac{2}{3}} \frac{\widehat{\delta\phi}}{m_P} + \frac{2}{n^2} \frac{\widehat{\delta\phi}^2}{m_P^2} - \frac{8\sqrt{2}}{3n^3\sqrt{3}} \frac{\widehat{\delta\phi}^3}{m_P^3} + \frac{20}{9n^4} \frac{\widehat{\delta\phi}^4}{m_P^4} - \dots \right) \dot{\widehat{\delta\phi}}^2. \quad (4.6b)$$

On the other hand, \widehat{V}_{IG0} in Eq. (3.11) can be expanded about $\langle \phi \rangle$ as follows

$$\widehat{V}_{IG0} = \frac{\lambda^2 m_P^2}{6c_{\mathcal{R}}^2} \widehat{\delta\phi}^2 \left(1 - \sqrt{\frac{2}{3}} \left(1 - \frac{1}{n} \right) \frac{\widehat{\delta\phi}}{m_P} + \left(\frac{7}{18} + \frac{1}{n} + \frac{11}{18n^2} \right) \frac{\widehat{\delta\phi}^2}{m_P^2} - \dots \right). \quad (4.6c)$$

From the expressions above, Eqs. (4.6b) and (4.6c), – which reduce to the ones presented in Ref. [11] for $n = 2$ – we can easily infer that $\Lambda_{UV} = m_P$ even for $n > 2$. For the model of Sec. 3.3 the expansion corresponding to \widehat{V}_{IG0} in Eq. (3.25b) reads

$$\widehat{V}_{IG0} = \frac{3\lambda^2 m_P^2}{100c_{\mathcal{R}}^{2-4/n}} \widehat{\delta\phi}^2 \left(1 - \sqrt{\frac{2}{3}} \left(\frac{7}{5} + \frac{5}{n} \right) \frac{\widehat{\delta\phi}}{m_P} + \left(\frac{391}{450} + \frac{77}{15n} + \frac{179}{18n^2} \right) \frac{\widehat{\delta\phi}^2}{m_P^2} + \dots \right). \quad (4.6d)$$

Since $2 - 4/n \geq 0$ for $n \geq 2$, we can conclude from Eqs. (4.6b) and (4.6d) that in this case also $\Lambda_{UV} = m_P$, in agreement with our findings in Sec. 4.1.

5 CONCLUSIONS

In this work we showed that a wide class of IG inflationary models can be naturally embedded in standard SUGRA. Namely, we considered a superpotential which realize easily the IG idea and can be uniquely determined by imposing two global symmetries – a continuous R and a discrete \mathbb{Z}_n symmetry – in conjunction with the requirement that inflation has to occur for subplanckian values of the inflaton. On the other hand, we adopted two forms of Kähler potentials, one corresponding to the Kähler manifold $SU(2, 1)/SU(2) \times U(1)_R \times \mathbb{Z}_n$, inspired by no-scale SUGRA, and one more generic. In both cases, the tachyonic instability, occurring along the direction of the accompanying non-inflaton field, can be remedied by considering terms up to the fourth order in the Kähler potential. Thanks to the underlying symmetries the inflaton, ϕ appears predominantly as ϕ^n in both the super- and Kähler potentials.

In the case of no-scale SUGRA, the inflaton is not mixed with the accompanying non-inflaton field in Kähler potential. As a consequence, the model predicts results identical to the non-SUSY case independently of the exponent n . In particular, we found $n_s \simeq 0.963$, $a_s \simeq -0.00068$ and $r \simeq 0.0038$, which are in excellent agreement with the current data, and $\hat{m}_{\delta\phi} = 3 \cdot 10^{13}$ GeV. Beyond no-scale SUGRA, all the possible terms up to the forth order in powers of the various fields are included in the Kähler potential. In this case, we can achieve n_s precisely equal to its central observationally favored value, mildly tuning the coefficient $k_{S\Phi}$. Furthermore, a weak dependance of the results on n arises with the lower n 's being more favored, since the required tuning on $k_{S\Phi}$ is softer. In both cases a n -dependent lower bound on $c_{\mathcal{R}}$ assists us to obtain inflation for subplanckian values of the inflaton, stabilizing thereby our proposal against possible corrections from higher order terms in Ω_H . Furthermore we showed that the one-loop radiative corrections remain subdominant during inflation and the corresponding effective theory is trustable up to m_P . Indeed, these models possess a built-in solution into long-standing naturalness problem [8, 10] which plagued similar models. As demonstrated both in the EF and the JF, this solution relies on the dynamical generation of m_P at the vacuum of the theory.

As a bottom line we could say that although no-scale SUGRA has been initially coined as a solution to the problem of SUSY breaking [13, 16] ensuring a vanishing cosmological constant, it is by now recognized – see also [11, 15, 18] – that it provides a flexible framework for inflationary model building. In fact, no-scale SUGRA is tailor-made for IG inflation since the predictive power of this inflationary model in more generic SUGRA incarnations is lost.

NOTE ADDED

When this work was under completion, results from BICEP2 experiment [26] were released according to which $r = 0.16^{+0.06}_{-0.05}$. If we combine this result with Eq. (3.6c) we find a simultaneously compatible region $0.06 \lesssim r \lesssim 0.11$ (at 95% c.l.) which, obviously, is not fulfilled by the models presented here, since the predicted r lies one order of magnitude lower – see Eq. (3.22) and comments below Eq. (3.31c). Although it is still premature to exclude any inflationary model with r lower than the above limit since the current data cannot definitively rule out other sources of gravitational waves – see e.g. Ref. [27] –, we would like to emphasize that our conclusion related to the predictive power of no-scale SUGRA, thanks to simplifications of Eq. (2.8), remains valid. Indeed, no-scale SUGRA can assist us to supersymmetrize models with more exotic frame functions – see e.g. Ref. [6, 28] – which could ensure consistency with the aforementioned r values. These constructions do not remain generically intact beyond no-scale SUGRA.

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